Lower critical field H_{c1} and barriers for vortex entry in $Bi_2Sr_2CaCu_2O_{8+\delta}$ crystals

M. Nideröst[1], R. Frassanito[1], M. Saalfrank[1], A.C. Mota[1], G. Blatter[2], V.N. Zavaritsky[3], T.W. Li[4], and P.H. Kes[4]
[1] *Laboratorium für Festkörperphysik, ETH Zürich, 8093 Zürich, Switzerland*
[2] *Theoretische Physik, ETH Zürich, 8093 Zürich, Switzerland*
[3] *Kapitza Institute for Physical Problems, 117334 Moscow, Russia*
[4] *Kamerlingh Onnes Laboratorium, Rijksuniversiteit Leiden, PO Box 9506, 2300 RA Leiden, The Netherlands*
(February 1, 2008)

The penetration field H_p of $Bi_2Sr_2CaCu_2O_{8+\delta}$ crystals is determined from magnetization curves for different field sweep rates dH/dt and temperatures. The obtained results are consistent with theoretical reports in the literature about vortex creep over surface and geometrical barriers. The frequently observed low-temperature upturn of H_p is shown to be related to metastable configurations due to barriers for vortex entry. Data of the true lower critical field H_{c1} are presented. The low-temperature dependence of H_{c1} is consistent with a superconducting state with nodes in the gap function. [PACS numbers: 74.25.Bt, 74.60.Ec, 74.60.Ge, 74.72.Hs]

High-temperature superconductors (HTSC's) are strongly type II and as such their $H-T$ phase diagram involves complete flux expulsion (Meissner state) at low fields and magnetic field penetration in the form of quantized flux-lines or vortices (mixed state) at higher fields. The onset of the mixed state, where vortex penetration becomes energetically favorable, is defined as the lower critical field $H_{c1} \propto 1/\lambda^2$ (here, λ is the London penetration depth). From the temperature dependence of H_{c1}, important information can be gained, particularly, regarding the symmetry of the superconducting state, since the appearance of gap nodes strongly modifies the T-dependence of the superfluid density and thereby the penetration depth $\lambda(T)$.

The experimental determination of H_{c1} has been a challenging and controversial problem since the beginning of HTSC research. Not only the values reported for the first flux penetration field H_p are scattered over a wide field range, but in strongly layered superconductors with H perpendicular to the planes, striking features have been observed, such as a marked upturn of H_p for temperatures $T \lesssim T_c/2$ [1–6]. The origin of the positive curvature of H_p at low temperatures has been the matter of various speculations. Different mechanisms have been proposed: Bean-Livingston surface barriers [1,3,4,6], bulk pinning [2,6], a modification of the character of the field penetration in layered structures [5], a low-temperature enhancement of the superconducting order parameter in the normal layers [7], etc.

For type II superconductors, there are at least two kinds of barriers which hinder the system from reaching a thermodynamic equilibrium state: (i) surface and geometrical barriers [8–10] governing vortex penetration into the superconductor, (ii) bulk pinning barriers governing vortex motion in the superconductor. For the determination of H_{c1}, the relevant barriers are those which govern vortex entry into the sample. Of particular interest in this context is the phenomenon of vortex creep over surface and geometrical barriers. From the theoretical point of view this subject has been discussed in Refs. [10] and [11], however, to our knowledge no systematic experimental study has been carried out so far.

In this letter we investigate the dependence of H_p on the magnetic field sweep rate dH/dt of isothermal magnetization curves for the strongly layered $Bi_2Sr_2CaCu_2O_{8+\delta}$ (Bi2212) material. This has been done for crystals with different cross-sections along the c-axis. For specimens with ellipsoidal cross-sections geometrical barriers can be neglected [10,12], so that the relevant barriers for vortex entry are expected to be surface barriers. Indeed, for our specimen with an ellipsoidal-like cross-section, the H_p vs. dH/dt dependence obtained at high sweep rates is well described by vortex creep over surface barriers [11]. However, for decreasing dH/dt rates, these barriers are observed to collapse. The subsequent saturation of H_p at lower rates indicates that the system is in equilibrium with respect to surface barriers. This is interpreted as strong evidence that we have reached the lower critical field H_{c1} in our measurements. The obtained low-temperature dependence of H_{c1} is in good agreement with recent microwave absorption measurements of the penetration depth λ [13,14]. On the other hand, for specimens with rectangular cross-sections, no significant vortex creep over the relevant barriers for vortex entry could be observed, in consistency with theoretical results for geometrical barriers [10]. The low-temperature upturns of H_p observed for specimens of either cross-section are then explained in terms of measurements where the system is out of equilibrium with respect to barriers for vortex entry.

The crystals under investigation have been grown with different techniques and have different shapes. The specimen with an ellipsoidal-like cross-section along the c-axis has been grown in a ZrO_2 crucible [15] and is approximately $1 \times 1.3 \times 0.05$ mm^3 in size. The specimens with rectangular cross-sections along the c-axis have been grown with the traveling solvent floating zone method [16] and are slightly larger in size. The experi-

ment is performed in a non-commercial SQUID magnetometer where the sample is stationary in the pick-up coil. The field H is supplied by a superconducting coil working in a non-persistent mode. For the magnetization curve measurements, the sample is zero field cooled from above T_c and stabilized at a fixed temperature (the residual field of the cryostat is $< 10\,\mathrm{mOe}$). Further, the field H is applied at a fixed rate dH/dt. For H_p, we select the field at which a deviation from Meissner shielding occurs (see inset of Fig. 3).

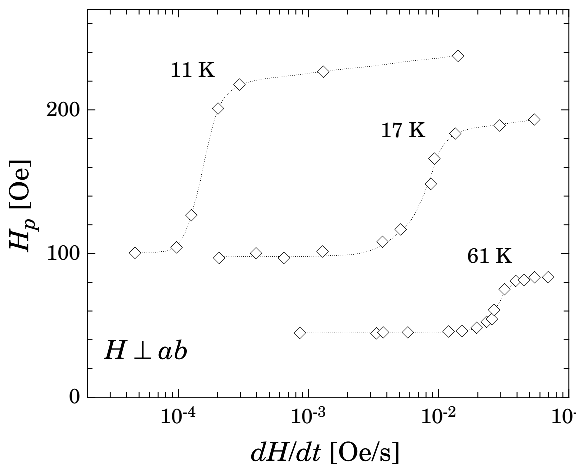

FIG. 1. For the specimen with an ellipsoidal-like cross-section, the first flux penetration field H_p vs. the applied magnetic field sweep rate dH/dt is displayed for different temperatures. The dotted lines are guides to the eyes. The data are scaled with the demagnetization factor $N = 0.96(\pm 15\%)$.

Fig. 1 shows the dependence of H_p on the field sweep rate dH/dt at three characteristic temperatures for the specimen with an ellipsoidal-like cross-section. A sharp step in the H_p vs. dH/dt curves is observed at all three temperatures. At high sweep rates, the curves display a finite slope which is most pronounced at $T = 11\,\mathrm{K}$ and $17\,\mathrm{K}$ and to a lesser degree at $T = 61\,\mathrm{K}$. These slopes are in good agreement with the theoretical results of Burlachkov *et al.* [11] for vortex creep over surface barriers (see the analysis below). Proceeding from the creep regime towards lower rates, the sharp drop signals a collapse of the surface barriers (or equivalently a heat pulse) which leads to a rapid flux entry, possibly in terms of flux jumps. Indeed, an avalanche type of flux penetration into a type II superconductor has been reported by Durán *et al.* [17]. On decreasing the sweep rates even further, the flux penetrates into the sample through the occurrence of rare events when the applied magnetic field reaches H_{c1}. We thus identify the saturated value of H_p with the lower critical field H_{c1} and use it to obtain information about the nature of the superconducting state.

Analogous measurements have been done on thin specimens with rectangular cross-sections. According to Zeldov *et al.* [10], for such specimens, geometrical barriers build up over a distance s from the sample edges, where s is the sample thickness. The energy required to overcome such a barrier is macroscopic $\sim \varepsilon_0 s$ so that vortex creep over geometrical barriers is expected to be very weak (here $\varepsilon_0 = (\Phi_0/4\pi\lambda)^2$ and Φ_0 is the unit flux). As a matter of fact, we were not able to detect significant vortex creep over geometrical barriers within the experimental time scale $10^{-5} \lesssim dH/dt \lesssim 10^{-1}$ Oe/s.

We proceed with a brief discussion of vortex creep over surface barriers. As shown by Burlachkov *et al.* [11], the surface barrier for pancake-vortices is given by $U \simeq \varepsilon_0 d \ln(0.76 H_c/H)$, where H_c is the thermodynamic critical field and d the interlayer distance. During the time t, thermal creep allows to overcome barriers of size $U(t) \simeq T\ln(t/t_0)$, where t_0 is a "microscopic" time scale [18]. Equating the two expressions one obtains the time dependence of the penetration field for the pancake-vortex regime,

$$H_p(t) \simeq H_c\,(t/t_0)^{-T/\varepsilon_0 d}\,. \qquad (1)$$

With the definition $T_0 = \varepsilon_0 d/\ln(t/t_0)$ Eq. (1) becomes $H_p \simeq H_c \exp(-T/T_0)$. At higher temperatures creep proceeds via half loop excitations of vortex-lines [11]. The creation of a half loop saddle configuration involves an energy $U \simeq \pi\varepsilon\varepsilon_0^2 c \ln^2(j_0/j)/2\Phi_0 j$, where j_0 is the depairing current density and $\varepsilon = (m/M)^{1/2} < 1$ the anisotropy parameter. With a similar analysis as above, the time dependence of H_p for the vortex-line regime takes the form

$$H_p(t) \approx H_c \frac{\pi}{2\sqrt{2}} \frac{\varepsilon\varepsilon_0\xi}{T\ln t/t_0} \ln^2\left(\frac{H_c}{H_p}\right), \qquad (2)$$

where ξ is the coherence length. According to Ref. [11], this expression gives a temperature dependence $H_p \propto (T_c - T)^{3/2}/T$. The crossover between the two regimes occurs when the size of the half loop excitation is of the order of the interlayer distance d. This is expected to occur at a temperature $T^* \approx T_0(T^*)\ln(d/\varepsilon\xi)$.

According to (1), creep of pancake-vortices over surface barriers results in a linear dependence of $\ln(H_c/H_p)$ on $\ln(t/t_0)$ with a slope $T/\varepsilon_0 d$. This dependence is given in Fig. 2(a) using the data at high cycling rates in Fig. 1. For $T = 11\,\mathrm{K}$ and $17\,\mathrm{K}$, the values of $T/\varepsilon_0 d$ obtained from the fits are in good agreement with the values estimated from the penetration depth $\lambda_{ab}(T)$ (see Table I).

TABLE I. $T/\varepsilon_0 d$ values calculated with the penetration depth $\lambda_{ab}(T)$ as obtained with formula $H_{c1} = (\Phi_0/4\pi\lambda_{ab}^2)\cdot\ln\kappa$ from the H_{c1} data in Fig. 3 and values of $T/\varepsilon_0 d$ obtained from the fits in Fig. 2(a).

	$T = 11\,\mathrm{K}$	$T = 17\,\mathrm{K}$
$T/\varepsilon_0 d$ (calculated)	1/34	1/21
$T/\varepsilon_0 d$ (from fit)	1/43 ±30%	1/22 ±30%

On the other hand, for $T = 61\,\mathrm{K}$ the above analysis for pancake-vortices is not satisfactory since the $T/\varepsilon_0 d$ value obtained from the fit is one order of magnitude smaller than the calculated value. According to (2), for creep of vortex-lines over surface barriers $(H_c/H_p) \cdot \ln^2(H_c/H_p)$ vs. $\ln(t/t_0)$ is linear with a slope $2\sqrt{2}T/\pi\varepsilon\varepsilon_0\xi$. For the $T = 61\,\mathrm{K}$ data at high cycling rates in Fig. 1, such a representation is given in Fig. 2(b). The slope $2\sqrt{2}T/\pi\varepsilon\varepsilon_0\xi$ obtained from the fit is 15 ± 3, in good agreement with the value 16, estimated with the help of the GL formula for the coherence length $\xi(T)$, with $\xi(0) \simeq 25\,\text{Å}$ and an anisotropy parameter $\varepsilon \simeq 1/100$. From these considerations we conclude that the behavior of H_p at high sweep rates is determined by creep of pancake-vortices (at low temperatures) and vortex-halfloops (at higher temperatures) over surface barriers. The estimated value of the activation barrier for vortex entry is $U \simeq 300\,\mathrm{K}$ for temperatures T between 10 and 20 K and $U \simeq 1200\,\mathrm{K}$ for $T \simeq 60\,\mathrm{K}$ (here we set the time origin $t = 0$ when the applied field H reaches H_{c1} and assume $t_0 \simeq 10^{-6}\,\mathrm{s}$).

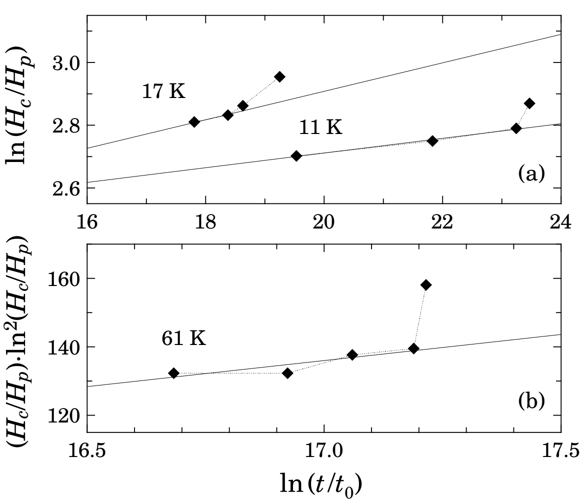

FIG. 2. (a) $T = 11\,\mathrm{K}$ and 17 K data of the upper plateau in Fig. 1 in a different representation. Here we choose the time origin $t = 0$ at $H = H_{c1}$. The critical field $H_c(T) \simeq H_c(0)\,(1 - T/T_c)$ is calculated with $\lambda_{ab}(0)$ as obtained from the H_{c1} data in Fig. 3 and the coherence length $\xi(0) \simeq 25\,\text{Å}$. We assumed $t_0 \simeq 10^{-6}\,\mathrm{s}$. (b) $T = 61\,\mathrm{K}$ data at high sweep rates in Fig. 1 in a convenient representation. The lines in Figs. 2(a) and (b) are fits according to Eqs. (1) and (2), respectively.

In Fig. 3 we make use of the curves in Fig. 1 to determine $H_p(T)$ with three different criteria: (i) H_p data ($\circ$) are taken at low sweep rates so as to guarantee that they lie in the regime where the system is in equilibrium with respect to surface barriers (these data represent $H_{c1}(T)$), (ii) H_p data ($*$) are taken at high rates so that they always lie in the creep regime, and (iii) the field H is swept at the constant rate $dH/dt = 8 \cdot 10^{-3}\,\mathrm{Oe/s}$ for all temperatures. In this case, the values of H_p (∇) lie on the H_{c1}-curve for $T \gtrsim 60\,\mathrm{K}$, they go through the step-like crossover for temperatures $30\,\mathrm{K} \lesssim T \lesssim 60\,\mathrm{K}$, and finally, for $T \lesssim 30\,\mathrm{K}$ they lie in the creep regime. From a comparison of the curves in Fig. 3, it follows that the low-temperature upturn of H_p results from measurements where the system is out of equilibrium with respect to surface barriers.

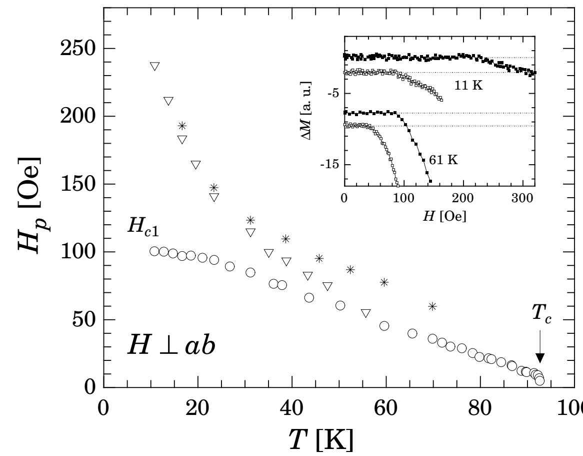

FIG. 3. Temperature dependence of the first flux penetration field H_p of the specimen with an ellipsoidal-like cross-section for different field sweep rates: ($\circ$) $dH/dt \leq 1 \cdot 10^{-4}\,\mathrm{Oe/s}$, ($*$) $dH/dt \geq 5 \cdot 10^{-2}\,\mathrm{Oe/s}$, and ($\nabla$) $dH/dt = 8 \cdot 10^{-3}\,\mathrm{Oe/s}$. The inset shows the deviation from the Meissner slope for $T = 11\,\mathrm{K}$ and 61 K. The upper curves ($\blacksquare$) are measured at a rate $dH/dt \geq 5 \cdot 10^{-2}\,\mathrm{Oe/s}$ and the lower ones ($\square$) at $dH/dt \leq 1 \cdot 10^{-4}\,\mathrm{Oe/s}$. The data are scaled with the demagnetization factor $N = 0.96(\pm 15\%)$.

As shown in Fig. 3, for increasing temperatures the T-dependence of H_p in the creep regime, ($*$) and (∇) for $T \lesssim 30\,\mathrm{K}$ and ($*$) for $T \gtrsim 30\,\mathrm{K}$, undergoes a crossover from an exponential behavior to a weak power-law. This is in agreement with the results presented above and indicates that the crossover from creep of pancake-vortices to creep of vortex-halfloops over surface barriers takes place at $T^* \approx 30 - 40\,\mathrm{K}$ (in consistency with the estimate $T^* \approx T_0(T^*)\ln(d/\varepsilon\xi) \approx 40\,\mathrm{K}$ obtained using the parameters for Bi2212 and $\ln t/t_0 \approx 20$).

Making use of the correct H_{c1} data for $H \perp ab$, ($\circ$) in Fig. 3, we have investigated the temperature dependence of the penetration depth λ_{ab} using the expression $H_{c1} = (\Phi_0/4\pi\lambda_{ab}^2) \cdot \ln\kappa$, (here we assume that $\ln\kappa$ is T-independent). For $11\,\mathrm{K} < T \lesssim 30\,\mathrm{K}$, we find a linear behavior $\lambda_{ab}(T) - \lambda_{ab}(0) \propto T$ with slope $d\lambda/dT \simeq 10\,\text{Å/K}$, where a linear extrapolation to the data gives a $T = 0$ penetration depth $\lambda_{ab}(0) \simeq 2700\,\text{Å}$. Our measurements thus confirm the linear low-temperature dependence of λ_{ab} as well as the slope $d\lambda/dT$, which have been recently determined with microwave absorption techniques on clean Bi2212 crystals [13,14]. A linear low-temperature dependence of λ_{ab}, as first observed by Hardy *et al.* [19] in YBCO crystals, is consistent with a

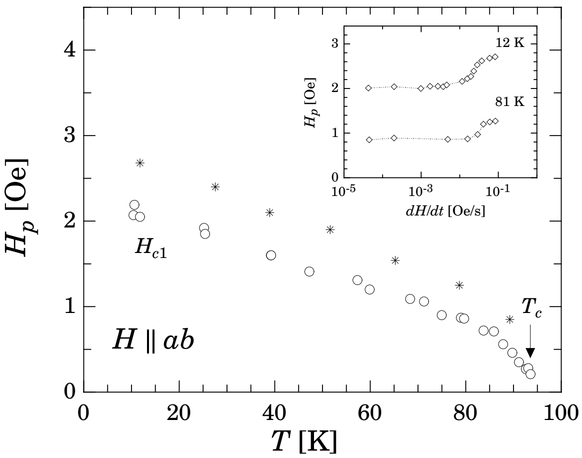

FIG. 4. Temperature dependence of H_p ($H \parallel ab$) for a specimen with rectangular cross-section and size $2\times3.9\times0.05$ mm^3 for different sweep rates: ($*$) $dH/dt \geq 6\cdot10^{-2}$ Oe/s and ($\circ$) $dH/dt \leq 6\cdot10^{-3}$ Oe/s. The inset shows the H_p vs. dH/dt dependence at $T = 12$ K and 81 K.

superconducting state with line nodes on the Fermi surface [20], e.g., a d-wave symmetry of the order parameter.

Similar investigations of H_p have been done for $H \parallel ab$. The misalignment angle between H and the ab-planes has been estimated to be smaller than 2° from measurements at magnetic fields in the range $H_{p\parallel} \ll H < H_{p\perp}$, (see Ref. [21]; here $H_{p\parallel}$ and $H_{p\perp}$ are the parallel and perpendicular penetration fields). For this configuration geometrical barriers are irrelevant. As shown in Fig. 4, for temperatures T between 10 and 70 K, $H_{c1}(T)$ has an approximately linear behavior. According to Ref. [22], for temperatures not so close to T_c, the correct description of H_{c1} in strongly layered materials is $H_{c1} = \Phi_0/(4\pi\lambda_{ab}\lambda_c)\cdot[\ln(\lambda_{ab}/d)+1.12]$. With this formula we calculated the penetration depth in c-direction λ_c using the previously determined λ_{ab} data. For $T \lesssim 40$ K, λ_c is approximately linear in T with a slope of $0.1\,\mu$m/K. The extrapolated $T = 0$ value is $\lambda_c(0) \simeq 15\,\mu$m. Fig. 4 further shows H_p data at high sweep rates ($*$). Contrary to the case for $H \perp ab$, no upturn of H_p is observed here for temperatures $T \lesssim T_c/2$. We attribute this difference to the absence of the strong pancake vortex creep regime at low temperatures for $H \parallel ab$. From the H_p vs. dH/dt dependence in the inset of Fig. 4, we obtain an activation barrier for vortex entry $U \simeq 200$ K at $T = 12$ K and $U \simeq 1300$ K at $T = 81$ K.

Close to T_c a downward bending of H_{c1} is observed for $H \perp ab$ as well as $H \parallel ab$ (see Figs. 3 and 4). This can be understood on the basis of an entropic downward renormalization of the vortex line free energy due to fluctuations of the order parameter around its mean-field form [23]. The decrease in the free energy $f_l = \varepsilon_l - Ts_l$ then leads to a drop in H_{c1} as T approaches T_c (here, ε_l is the line energy of the vortex excitation and s_l is the line entropy).

To conclude, isothermal magnetization curves were measured on Bi2212 crystals for configurations of negligible geometrical barriers. At very low sweep rates we found saturated values of the penetration field H_p which we interpret as the true lower critical field H_{c1}. The low-temperature dependence of H_{c1} is consistent with recent magnetic penetration depth measurements [13,14] suggesting a superconducting state with nodes in the gap function. Based on the results obtained for Bi2212, we argue that the frequently reported low-temperature upturns of H_p in strongly layered superconductors with H perpendicular to the planes [1–6] can be explained in terms of measurements where the system is out of equilibrium with respect to barriers for vortex penetration.

It is a great pleasure to acknowledge K. Aupke, C. de Morais-Smith, V.B. Geshkenbein, A. Suter, T. Teruzzi, and P. Visani for fruitful discussions and helpful contributions. This work has been supported by the Swiss National Science Foundation.

[1] V. N. Kopylov *et al.*, Physica C **170**, 291 (1990).
[2] N. V. Zavaritsky and V. N. Zavaritsky, Zh. Eksp. Teor. Fiz. **53**, 212 (1991), [JETP Lett. **53**, 226, (1991)].
[3] A.C. Mota *et al.*, Physica C **185–189**, 343 (1991).
[4] N. Chikumoto *et al.*, Physica C **185–189**, 1835 (1991).
[5] V. V. Metlushko *et al.*, Phys. Rev. B **47**, 8212 (1993).
[6] E. Zeldov *et al.*, Europhys. Lett. **30**, 367 (1995).
[7] T. Koyama and N. Takezawa and M. Tachiki, Physica C **168**, 69 (1990).
[8] C. P. Bean and J. D. Livingston, Phys. Rev. Lett. **12**, 14 (1964).
[9] M. V. Indenbom *et al.*, Physica C **222**, 203 (1994).
[10] E. Zeldov *et al.*, Phys. Rev. Lett. **73**, 1428 (1994).
[11] L. Burlachkov *et al.*, Phys. Rev. B **50**, 16770 (1994).
[12] D. Majer *et al.*, Phys. Rev. Lett. **75**, 1166 (1995).
[13] T. Jacobs *et al.*, Phys. Rev. Lett. **75**, 4516 (1995).
[14] S.-F. Lee *et al.*, Phys. Rev. Lett. **77**, 735 (1996).
[15] N. V. Zavaritsky, A. V. Samoilov, and A. A. Yurgens, Physica C **169**, 174 (1990).
[16] T. W. Li *et al.*, J. Crystal Growth **135**, 481 (1994).
[17] C. A. Durán *et al.*, Phys. Rev. B **52**, 75 (1995).
[18] G. Blatter *et al.*, Rev. Mod. Phys. **66**, 1125 (1994).
[19] W. N. Hardy *et al.*, Phys. Rev. Lett. **70**, 3999 (1993).
[20] J. F. Annett, N. Goldenfeld, and S. R. Renn, Phys. Rev. B **43**, 2770 (1991).
[21] N. Nakamura, G. D. Gu, and N. Koshizuka, Phys. Rev. Lett. **71**, 915 (1993).
[22] J. R. Clem, M. W. Coffey, and Z. Hao, Phys. Rev. B **44**, 2732 (1991).
[23] G. Blatter, B. Ivlev, and H. Nordborg, Phys. Rev. B **48**, 10448 (1993).